

KANSAS 94-5-22
OKHEP-94-02

Evidence for Gluon Energy Loss as the Mechanism for Heavy Quarkonium Suppression in pA Collisions

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Abstract

We study the energy and nuclear A dependence of the hadronic production of heavy quarkonia. We review theoretical ideas which have been put forward, seeking a consistent global picture reconciling the large effects in quarkonia with the small nuclear effects observed in continuum Drell Yan production. The data indicates that shadowing or leading twist modifications of parton distributions can be ruled out as explanations, leaving higher twist energy loss. From general principles the maximum allowed energy loss of partons traversing the nuclear medium can be related to the parton transverse momenta. We then show that the experimental data on nuclear suppression of charm- and bottom- onium for large x_F is consistent with this effect: using the observed transverse momenta to bound the x_F dependence in an almost model independent manner generates a relation that practically reproduces the data. Several prediction are discussed; the dependence on x_F as $x_F \rightarrow 1$, and large and small k_T^2 cuts, can be used to discriminate between quark and gluon induced effects.

1. INTRODUCTION

Propagation of quarks, gluons and hadrons through nuclear matter is currently a subject of intense interest. It is expected that the study will teach us much about the interplay between perturbative and non-perturbative QCD. An important experimental discovery by the Fermilab E772 experiment [1] is the suppression of charmonium and bottomonium production in pA collisions in comparison to the production rate in pp collisions. A similar effect had also been previously reported by Badier et al [2] and Katsanevas et al. [3]. The data strongly contradicts a widespread theoretical expectation that at high energies the nuclear medium should have negligible effect on heavy quarkonium production. The observed suppression has direct implications for the use of onium production as a signal for quark-gluon plasma formation in heavy ion collisions. It has generated much controversy [4-7], and raised the possibility that the observations represent a serious challenge to theory.

A common theoretical prejudice that suggests negligible nuclear suppression is the following. One can argue that the characteristic separation of the charmed quark anti-quark pair produced by a partonic interaction is of the order of $1/m_c$, which is equal to $(1.5 \text{ GeV})^{-1}$. Invoking a geometrical cross section in the spirit of color transparency, the attenuation cross section of the charm pair might be as small as the order of $(1.5 \text{ GeV})^{-2}$, which is about 0.2 mb. With this cross section, a typical survival factor of 0.97 is obtained for a large nucleus of diameter 10 fm. This is a very small effect in comparison to 30 to 60 % suppression seen in the data.

Of course, such a method of estimating the nuclear effects applies (at best) to the propagation of a well-localized, relativistic, color-singlet charm pair. But invoking color transparency for the actual production is probably not correct, since the kinematics of the events are highly inelastic, and lack the usual conditions of exclusivity that color transparency arguments should assume. Color transparency is not expected to occur if the coherence of a system is broken, for example in the case when uncontrolled inelastic color flows are summed over in a semi-inclusive production. Moreover, in the initial state the gluons can lose energy by interacting with the nuclear medium, and

there also is the likelihood that the charmed pair is produced temporarily in the color octet (rather than singlet) state. This does not mean that the suppression is of the expected size: indeed, the huge magnitude of the effect has been quite mysterious.

Besides the outstanding puzzle of quarkonium suppression, the general problem of parton propagation in the nucleus is becoming rather important. There is great interest in understanding the role of quantum mechanical coherence of QCD interactions in the nuclear medium. Depending on the experimental circumstances, the same arguments leading to coherent suppression of energy loss in high energy QED interactions - the so called LPM [8-10] effect - can be applied, leading to many interesting predictions. In the LPM effect, a concept of “formation time” τ_{form} of quanta during a high energy interaction is considered. The order of magnitude of the formation time for a quantum of mass m , carrying energy E , and transverse momentum k_T with respect to its progenitor is

$$\tau_{\text{form}} \approx E/(k_T^2 + m^2) \quad (1)$$

The LPM effect in QED suppresses bremsstrahlung due to destructive interference between emissions occurring over an interaction time $\tau_{\text{int}} \leq \tau_{\text{form}}$. A recent paper by Gyulassy and Wang [11] studies the effects of multiple scattering in perturbative QCD. Although the details of coherence are considerably more complicated due to the non-Abelian color algebra, the basic features of the LPM effect are not changed by this study.

The concept of the formation time leads to a complementary concept of formation rate Γ_{form} . The formation rate is simply the inverse of the formation time,

$$\Gamma_{\text{form}} \approx (k_T^2 + m^2)/E \quad (2)$$

Examined in this way, there is an interesting possibility that the coherent parton formation rate could be “tuned” by selecting signals with various masses, energies and transverse momenta. Is it possible, for example, to increase particle formation by selecting events with larger k_T^2 ? As we will show below, the trends in the data indicate possible observation of a dramatic “anti-LPM” enhancement of parton emission due to increased formation rate

associated with large k_T^2 . The possibility of tuning the formation rate leads naturally to a number of interesting signals which can be experimentally tested.

2. REVIEW

In view of these points (and the historical difficulty of understanding charm production in general), we will approach the problem of onium suppression in a methodical manner, and try to limit our assumptions carefully:

i) Electroproduction experiments have convincingly shown that the process of parton fragmentation is negligibly affected by nuclei. The E772 experiment itself also found that continuum Drell-Yan muon pair production shows no significant nuclear effect. The data for continuum dimuon $A^{-1}d\sigma/dQ^2$, for example, shows little nuclear effect in the Q^2 regions both above and below the suppressed onium regions. These observations are consistent with conventional factorization. They seem to significantly limit the amount of energy loss one can attribute phenomenologically to parton propagation inside the nuclear medium.

However, we note that both the bulk of fragmentation and the Drell Yan continuum are mainly probes of the quark-antiquark channel. They say little about production via gluons. The nature of gluon and heavy quark propagation in nuclei is not clear.

ii) A very important clue to the charmonium data is given by the fact that the data's x_2 dependence does not scale. Even if one creates an ad-hoc gluon or quark-antiquark distribution for one experiment on charm, it does not reproduce the data for bottom. Furthermore the data does not scale with x_2 as we change the incident beam momentum from 200 to 800 GeV [12]. This eliminates the possibility of gluon shadowing or more general EMC-type effects on the parton distributions as a dominant mechanism. Quite recently, Banesh, Qiu and Vary [7] (BQV) claim that the onium production at large x_F is dominated by quark annihilation, and that shadowing in the small- x_2 quark distributions accounts for the dominant part of the suppression. Of course, shadowing makes a contribution, but the magnitude of small x shadowing observed in deeply inelastic scattering is not nearly large enough

to account for the onium suppression. Postponing a detailed discussion to Section 5, we do not believe that the BQV analysis actually support the claim. The same objection based on factorization must be raised here: the onium production data shows that factorization does not hold, ruling out this mechanism as a model for the full effect.

- iii) Since factorization is a leading twist prediction for this kind of semi-inclusive production, and the same sort of experiment has confirmed its use and understanding elsewhere, the suppression effect must be a higher twist one. The bizarre thing is that it so drastically affects the onium production.
- iv) The quarkonium suppression seen in the E772 data can be questioned as to whether it might be an instrumental effect. In fact, certain regions of the data's transverse momentum spectra have not been reported due to questions about the acceptance [13]. We re-examined previous data of Badier et al [2] on J/ψ production on nuclear targets in studying this. We find that rather than contradicting the E772 experiment, the Badier et al data seems to confirm the trend. A signal for suppression of onium is real and more than ten years old.

We now summarize some theoretical ideas, which we consider as "spare parts" that might be assembled in a new order to understand the puzzle:

- a) It can be argued that heavy quarkonium, as a non-relativistic bound state, might have a large nuclear matter interaction cross section that is set by the full onium size and binding energy rather than the mass. The J/ψ bound state diameter, for example, is 3 to 5 times the charmed quark compton wavelength, leading to an estimated cross section as large as 2 to 5 mb. A major problem with this proposal is that the time scale for formation of a bound state onium is very long compared to the time scale for crossing the nucleus. This is aggravated by the fact that the Lorentz boost stretches any effective time scale enormously, so that a J/ψ does not become a J/ψ until it is more than 100 Fm away! We do not pursue this idea further.
- b) There remains a realistic possibility that interaction of incoming and outgoing colored partons could cause them to lose energy. Gavin and Milana (GM) [5] observed that even a small shift in the x_F value of the incoming partons, assumed by them to be gluons, could lead to a numerically large

change in $d\sigma/dx_F$ because of the rapid variation with x of the gluon distribution functions.

For example, supposing the gluon distribution to go like $(1-x)^5$, then a gluon taken from the gluon distribution at $x + \delta x$ is less likely to contribute by the ratio $(1-x-\delta x)^5/(1-x)^5 \approx 1 - 5\delta x/(1-x)$. Even if δx is small, this creates a strong kinematic suppression as x approaches unity. For reference we will call this the GM mechanism. This effect was assumed to occur for both initial and final state propagation, with onium production dominated by color octet components.

To get a large enough shift in the x value, GM also proposed a rule for energy loss which is of the “higher twist” type; their proposal for the energy loss is $\Delta E/E = cx_1 A^{1/3}/Q^2$, where c is a color dependent constant that could be adjusted to fit to the data. At first the higher twist character of the GM rule seemed to put it into a phenomenological limbo of the uncalculable, a thing which could be neither verified nor disproven with current theoretical knowledge.

c) This approach was countered by Brodsky and Hoyer (BH) [6] who argued that the dependence on energy of the GM formula violates general principles. BH went on to claim that there exists an upper bound on the energy loss for a parton propagating through nuclei. This upper bound, in the spirit of LPM, is obtained from rather general considerations; the BH rule is $\Delta x < k_T^2 L_A / 2E$, where k_T is the transverse momentum change in the collisions giving the energy loss and L_A is the target length. This rule also is higher twist, but the uncalculable higher twist mechanism of Gavin and Milana has a limit that contradicts this relation, allowing ΔE to go like E at fixed Q^2 . The contradiction between the two formulas is numerically important. Using their own bound for the energy loss and a value for $k_T = 300$ MeV, Brodsky and Hoyer dismissed the Gavin and Milana proposal as insufficient to explain the data.

3. ASSEMBLING THEORETICAL SPARE PARTS

Considering these ideas, let us first observe that the GM mechanism is (a) quite reasonable and (b) logically independent of the model used for the energy loss. A first task, then, is to determine whether some energy loss, however it occurs, can explain the data. This has already been answered by GM who were able to fit the data. There is no reason, then, to rule out gluon energy loss as a mechanism, but we agree that one should have a consistent framework to represent it.

A second observation is that the BH expression for energy loss can be viewed a testable hypothesis, that is, as a model. The rule is claimed to be an upper bound, so any observed energy losses must lie inside an envelope of values it specifies. This prompts us to examine the experimental data, using the GM hypothesis, and test for the general mechanism of energy loss by seeing whether or not the data obeys the claimed bound. Our goal here is to check the energy loss proposal without getting snarled into model dependence of the energy loss formula.

Third, given the fact that Nature tends to dissipate energy rather maximally, we can try saturating the bound and examining whether we obtain a prediction close to the behavior shown in the data. Actually, by examining the bound more closely, we find it needs modification from mass effects and also by an unknown dimensionless factor. We consider our additions to be modest corrections. The trends in the data are not strongly dependent on the unknown prefactor.

The arguments leading to the BH bound, its modifications, and its application along with the GM mechanism are reviewed in the next section. We find that we can test for energy loss by examining the detailed x_F distribution using the k_T distribution as an input. This is a much more powerful test than simply looking at overall production rates. The procedure works quite well in the upsilon meson case where the k_T distribution has been measured. We find that the suppression seen in the x_F distribution is well within the limits imposed by the bound on energy loss. Even more importantly, the dependence of the data on x_F and A actually tend to parallel the bound. We consider the agreement of a general bound and a hitherto unnoticed pattern

in the data itself to be practically model independent evidence that the basic culprit in nuclear quarkonium suppression is gluon energy loss.

While the bottom quark case checks well, a definitive application of this result to E772 charmonium production is not yet possible because of problems due to experimental acceptance [13]. The transverse momentum distribution for the charmed case is known only for k_T less than about 2 GeV. We can nevertheless apply our formalism in reverse and use the experimentally observed x_F distribution to put a lower limit on the k_T distribution of charmonium. This prediction can be tested in future experiments. Our results show that the lower limit on the k_T distribution is roughly the same as the corresponding observed distribution for the case of bottomonium.

3.1 THE ROLE OF GLUONS

Although we will present evidence that the GM mechanism is at work, this connection seems, at first sight, not specific to the production mechanism. There is a problem due to the generality of the BH bound. The BH bound assumes too little; one might say it tells us that we are seeing the uncertainty principle at work, a fact of limited usefulness. One knows energy was lost but does not identify the underlying subprocess with this mechanism alone.

The surprise of the onium suppression is contained in the large magnitude of the effect. One of our main points is that the quarkonium suppression (which has gotten so much attention) is directly related to the transverse momentum in the data (which has gotten very little attention). The average transverse momentum squared $\langle k_T^2 \rangle_A$ is rather large;¹ for Tungsten the $\langle k_T^2 \rangle_A$ for bottomonium is greater than 2 GeV² rather than 0.1 GeV². As a function of A , the integrated Υ_{1s} data is well described by [1,14]

$$\langle k_T^2 \rangle_A = [0.16(A^{1/3} - 2^{1/3}) + 2.59] \text{ GeV}^2 . \quad (3)$$

The same experiment found the A dependent part of the Drell Yan continuum $\langle k_T^2 \rangle_A$ to be about 10 times smaller. As we will show, much can be

¹A similar observation on the largeness of nuclear interaction induced transverse momentum in dijet production has been made by T. Fields and M. Corcoran (to appear in *Proceedings of EPS Conference*, Marseille 1993). We thank Tom Fields for informing us of this.

predicted simply knowing the $\langle k_T^2 \rangle_A$.

What, then, is causing the rapid variation with A of the quarkonium $\langle k_T^2 \rangle_A$? Variation of mean transverse momentum with A has also been seen in other experiments studying nuclear dependence of dijet and dihadron production [15,16]. Physically this effect can arise from multiple scattering in the nuclear medium [17]. The data's $A^{1/3}$ dependence indicates a target length effect, consistent with the energy loss mechanism. The same data indicates that the charged partons contributing to Drell Yan are acting different from the partons making the heavy quarkonia. Drell Yan precludes any strong effects for the quarks, then! This suggests that gluons, which give dominant contribution to heavy quarkonia production, are scattering more in the transverse direction and are losing more energy than quarks. An alternate possibility, which is not strictly ruled out by the current data, is that the final state interactions of the heavy quarks are responsible for the large nuclear effects. However, experience with QED bremsstrahlung and perturbative QCD, where a quark mass acts to substantially cut-off massless vector emissions, makes it hard to believe that heavy quarks could lose energy so much faster than light quarks! In Section 5 we show that analysis of data at large x_F will also be able to discriminate clearly between these two possibilities.

We believe that further data and not theoretical arguments will play the most important role in finally pinning down all the unknowns.

For better or worse, the dynamics of gluon channels producing onium have never been clear, and a number of issues, including “intrinsic charm”, have been created to address the problem. One cannot assign a perturbative overall normalization to the gluon channels safely. It is not known, e.g. what fraction of the time a color octet $q\bar{q}$ state is produced rather than a color singlet. For this reason, we have arranged our calculations so as not to base them on normalization factors. We will postpone to Section 5 a discussion of the interplay of subdominant quark channels with the gluon channels.

3.2 ENERGY LOSS RATES

We first review and expand on the argument of Ref. [6] to obtain an

expression for an upper bound on energy loss. Assume a parton propagating through the nuclear medium in the $+z$ direction with energy $E = x_1 E_p$, where E_p denotes the energy of the incoming proton. This parton loses an energy $x_g E$ by emitting a gluon in the presence of a source as shown in Fig. 1. This, and any secondary scattering, must occur in the volume of a nucleus. The finite size of the nucleus introduces a distance scale Δx , which will be used with the uncertainty principle to bound the energy loss.

Components of momentum q exchanged with the source can be related to the source spatial size Δx by $\Delta x \Delta q \geq 1$. The “ $-$ ” component of the exchanged momentum is easily probed. Letting the four momentum squared of the final state be given by M^2 , then $M^2 = q^2 + 2p \cdot q$. We solve for $q^- \approx \Delta q^- \approx M^2/(2p^+) - q^2/(2p^+)$, where $p^+ = (E + p_z)/\sqrt{2}$

Assume for definiteness now that we have two identical particles of mass m in the final state. Their momenta in $(+, T, -)$ notation can be listed as

$$\begin{aligned} p_1 &= [x_g E, \bar{k}_T, \bar{k}_T^2/(2x_g E)] \\ p_2 &= [(1 - x_g) E, -\bar{k}_T, \bar{k}_T^2/(2(1 - x_g) E)] \end{aligned} \quad (4)$$

where $\bar{k}_T^2 = k_T^2 + m^2$ is the “transverse mass”. The final state mass is

$$M^2 = \bar{k}_T^2 \left(\frac{x_g}{1 - x_g} + \frac{1 - x_g}{x_g} + 2 \right) . \quad (5)$$

Assuming $\bar{k}_T^2 \ll M^2$, then the kinematics forces either $x_g \ll 1$ or $(1 - x_g) \ll 1$. For gluons emitting gluons the two situations are physically identical. Thus it is a good approximation to take

$$M^2 \approx \frac{\bar{k}_T^2 + m^2}{x_g} \quad (6)$$

Combining the above with $\Delta q^- = (M^2 - q^2)/(2p^+)$, we have

$$1/\Delta q^- \approx 2p^+ x_g / (\bar{k}_T^2 + m^2 - x_g q^2) , \quad (7)$$

This can be combined with the null plane uncertainty principle $\Delta x^+ \Delta q^- \geq 1$ to give

$$2p^+x_g/(k_T^2 + m^2 - x_g q^2) < \Delta x^+ , \quad (8)$$

The meaning of Δx^+ is the light-cone “time” within which the events occur. This is clarified by a space-time cartoon (Fig 2). From the figure, due to the finite size of the nucleus, the Δx^+ for events causally propagating along the future light cone obeys $\Delta x^+ < \sqrt{2}L_A$, where L_A is the rest frame length of the nucleus. Note that non-relativistic propagation would allow Δx^+ to become indefinitely large if the parton stops inside the nucleus. This possibility is irrelevant to the discussion, although it would weaken the bound.

BH did not include the effects of m^2 and q^2 . Ordinarily such terms are negligible in comparison with large momentum transfers. However in this case we see that they are not necessarily negligible in comparison with k_T^2 . Setting q^2 and m^2 to zero in (8), we obtain the result of BH

$$x_g \leq L_A k_T^2 / (2E) \quad (9)$$

using $p^+ = \sqrt{2}E$. The energy loss is then given by $\Delta E_{\text{parton}} \approx x_g E = x_g x_1 E_p$, leading to $\Delta x_1 = x_g x_1 \leq L_A k_T^2 / 2E_p$, as found by BH.

Restoring the dependence on m^2 and q^2 , we find:

$$x_g \leq \frac{(L_A/2E)(k_T^2 + m^2)}{1 - |q^2|L_A/(2E)} . \quad (10)$$

At this point, if we know $|q^2|L_A/2E \ll 1$, then we have a BH bound with only the modification of the “transverse mass”. For our further discussion we will assume this limit.

3.3 BACKUP

Are there any important loopholes in the bound? We note the following:

- i) The neglect of a final state mass for the parton is dangerous. In the BH calculation, using k_T values around the traditional values of 300 MeV, then $k_T^2 = 0.1 \text{ GeV}^2$, which is much smaller than any value one would use for an effective m^2 , even considering “massless partons”. The bound is quite sensitive to this. We will use k_T^2 values obtained from data. This is an important detail for application of the bound, but not for its concept.

ii) The bound, while invoking general kinematics and quantum principles, nevertheless depends on the dynamics which assumed the production of *two particles per vertex* in the final state. It is similar in spirit to electrodynamics, where the LPM effect [8-10] provides the classic prototype for coherent suppression of energy loss in a finite density target.

To see this, we present a “back of the envelop” discussion of the physics of LPM, abstracted from Feinberg and Pomeranchuk [10]. The kinematics assume nearly forward scattering of an electron (mass = m_e , energy = E_e) in a classical medium, with emission of a bremmstrahlung photon carrying energy ω . All transverse momentum are assumed to be very small - this is an important point. For forward scattering, one finds a spatial momentum transfer $q_{||}$ given by

$$q_{||} = \sqrt{E_e^2 - m_e^2} - \sqrt{(E_e - \omega)^2 - m_e^2} - \omega \quad (11)$$

so that

$$\Delta q_{||} \approx m_e^2 \omega / (2E_e^2) \quad (12)$$

This momentum transfer determines a coordinate space region of longitudinal length $r \sim 1/\Delta q_{||} \sim 2E_e^2/(m_e^2 \omega)$. Suppose that the length is so large that the electron undergoes several multiple scatterings. From multiple scattering theory, the electron does a random walk with mean square scattering angle

$$\theta_s^2 \approx \left(\frac{E_s}{E} \right)^2 \frac{r}{L} \quad (13)$$

where E_s is an energy scale and L a radiation length. Note that θ_s^2 is proportional to the distance travelled r . The LPM argument, which is semiclassical, observes that coherence can be retained in bremmstrahlung emission over a cone given by an angle set by the boost parameter $\gamma = E/m$. Most photons are emitted with angle $\theta_\gamma \leq 1/\gamma$ relative to the moving source. Coherence over the emission becomes crucial if the multiple scattering angle θ_s is bigger than θ_γ . This condition is

$$\left(\frac{E_s}{E_e} \right)^2 \frac{r}{L} \geq \left(\frac{m}{E_e} \right)^2, \quad (14)$$

which upon inserting $r \sim 2E_e^2/m_e^2 \omega$ becomes

$$\omega \leq \frac{2E_e^2}{m_e^2} \frac{E_s^2}{m_e^2 L} \quad (15)$$

Emissions satisfying this bound are suppressed by destructive interference.

The emphasis in the LPM analysis is on a region where the photon formation rate is small compared to the collision rate. But then it follows that the process is exquisitely sensitive to the scales setting the rate $\Gamma_{\text{form}} = k_T^2/2\omega$. By adjusting k_T^2 we can evidently scan across a range of formation times, and turn the LPM suppression on or off. We will apply this observation to the quarkonium suppression in the next section.

Continuing, the dynamics of QCD has features which are different from QED. In QCD an incoming gluon can split into three gluons at a single interaction due to the perturbative 4-gluon vertex (Fig. (3)). This upsets the QED-based argument. We have worked through the kinematics of three particles in the final state, verifying that some regions reproduce the formation time argument while other regions exist which do not.

For definiteness consider Fig. (3), in which an incoming gluon splits into three gluons carrying momentum fractions x_1 , x_2 and $x_3 = 1 - x_1 - x_2$. The momenta are given by

$$\begin{aligned} P_1 &= \left(x_1 P, \quad k_{T,1}, \quad \frac{k_{T,1}^2}{2x_1 P} \right) \\ P_2 &= \left(x_2 P, \quad k_{T,2}, \quad \frac{k_{T,2}^2}{2x_2 P} \right) \\ P_3 &= \left(x_3 P, \quad k_{T,3}, \quad \frac{k_{T,3}^2}{2x_3 P} \right) \end{aligned}$$

in $(+, T, -)$ notation, assuming q^2 and $q_T^2 \rightarrow 0$, and using massless gluons. The q^- momentum is simply given by the sum of the p_i^- :

$$q^- = \frac{k_{T,1}^2}{2x_1 P} + \frac{k_{T,2}^2}{2x_2 P} + \frac{k_{T,3}^2}{2x_3 P}$$

Applying $q^- > 1/L$, where L is some interaction length, we can bound this sum.

The “formation time” is $\tau_{\text{form}} = 1/q^-$. We see that its inverse is the sum of three inverse formation times,

$$1/\tau_{\text{form}} = 1/\tau_1 + 1/\tau_2 + 1/\tau_3 \tag{16}$$

where each formation time (up to trivial factors) equals the usual definition $1/\Gamma_i = \tau_i \sim 2x_i P/k_{Ti}^2$. The interpretation of (16) is of course simpler in terms of rates: the total formation rate Γ_{TOT} is the sum of the three individual formation rates.

From the uncertainty principle the formation of the final state requires

$$\Gamma_{TOT} L \geq 1$$

or

$$L_A \geq \frac{1}{\sum_i \Gamma_i} \quad (17)$$

The finite value of L_A means that not all the formation rates can be too small or else they will destructively interfere.

Unlike the case of 2-body formation - where the creation of one particle implies the other - when we have a 4-point interaction, there is more than one independent formation rate. In the relation (17) the biggest formation rate wins. That is, in the region

$$\frac{k_{T,1}^2}{2x_1 P} \gg \frac{k_{T,2}^2}{2x_2 P} + \frac{k_{T,3}^2}{2x_3 P}$$

then the creation of parton “1” dominates the issue of coherent formation. The formation of partons “2” and “3”, while occurring less rapidly, is triggered by the formation of parton “1” at the 4-point vertex, *but they are not separately resolved at this point*.

Physically, here is what happens. An incoming gluon can be disrupted by the source into a small x gluon (say), while two larger- x comoving gluons are simultaneously created. The uncertainty principle applies to the smallest- x gluon which cannot be produced too slowly. But understanding the coherence and quantum mechanical resolution of the smallest- x gluon says *nothing about resolving two other gluons into separate components*. A subsequent hard collision (or similar independent time scale) is needed to resolve them.

Naturally we have a probe of a gluon’s momentum when a heavy quark is produced. Suppose, between x_2 and x_3 , we detect a quark carrying $x_2 \approx x_F$; what fixes the value of x_3 ? It is fixed by detailed dynamics, not general principles. This situation is unprecedented; it indicates the possibility of

energy losses not following the rules of QED. We conclude that the usual QED-LPM arguments are inadequate for a quantitative analysis of energy loss in QCD.

This is a real loophole; how big are its effects? The four gluon vertex is higher order in perturbation theory, being of order g^2 , but cannot be negligible because it is necessary for the gauge invariance of the theory.² Moreover, the problematic integration region is comparable or larger in size to the region being discussed in the QED formation time arguments. Yet our analysis of the dimensionless 4-gluon vertex emission does not introduce any new “large” scales beyond L_A , k_T^2 , and $1/E$ already present. By dimensional analysis, then, something like the LPM analysis should survive after detailed integrations and combinatorics over dimensionless quantities are evaluated. We will accept it for this study. We believe that further work could show that the bound might be multiplied by a dimensionless factor we estimate to be a few units.

4. APPLICATIONS

We now turn to applying the energy loss bounds as practical tools. The average value of k_T^2 in our formula represents the transverse momentum caused by scattering. We assume that this is the difference of intrinsic transverse momentum $\langle k_T^2 \rangle_{\text{int}}$, of the order of 0.5 - 1.0 GeV², and the observed value. A more crisp definition can be given for “leading twist” reactions but does not exist for power suppressed processes. To proceed orderly we separate the initial state energy losses of gluons from final state ones of heavy quarks. We present bounds based on initial state gluon energy losses only. We will show that we produce a trend that is strikingly consistent with the data. We discuss separating initial from final state effects in Section 5.

First we examine an analytic estimate of the energy loss. We assume gluons are bremmed off with a distribution

$$\frac{dN}{dx_g} = f(x_g) ; \quad f(x_g) = 0 , \quad x_g > x_{\max} .$$

²It has always been worrisome that the 4-gluon vertex has produced few qualitative effects in high energy phenomenology.

where x_{\max} is given by the bound (9). Since we are in quasi-nonperturbative region the details of f_g are unknown. So long as $f(x_g)$ is peaked at small x_g but regular at $x_g = 0$ the details turn out not to matter much.

The effective gluon distribution $\bar{G}(x)$ due to the shift from energy loss is given by

$$\bar{G}(x) = \int_0^1 dx_g f(x_g) (1 - x - x_g)^5$$

Using for example $f(x_g) = 1/x_g$, $x_{\min} < x_g < x_{\max}$, then we are interested in the limit $x_{\max} \ll 1$. A series expansion can be obtained by integrating by parts

$$\begin{aligned} \int_{x_{\min}}^{x_{\max}} dx_g \left[\frac{d}{dx_g} \ln x_g \right] (1 - x - x_g)^5 &= \ln x_g (1 - x - x_g)^5 \Big|_{x_{\min}}^{x_{\max}} \\ &\quad + 5 \int_{x_{\min}}^{x_{\max}} dx_g \ln(x_g) (1 - x - x_g)^4 . \end{aligned}$$

The second term can be integrated by parts again leading to an asymptotic series of any order desired. The first term is approximately

$$\bar{G} \sim \ln(x_{\max}/x_{\min}) (1 - x)^5 \left[1 - \frac{5x_{\max}}{1 - x} + \dots \right]$$

to first order in x_{\max} and dropping terms proportional to x_{\min} . The logarithm is slowly varying in both x_{\max} and the infrared cutoff x_{\min} and will be ignored. One sees the GM mechanism emerging in the power series expansion: the effects of a small x_{\max} get big as $x \rightarrow 1$.

Now suppose that \bar{G} is used in a calculation of cross section, namely

$$\frac{d\sigma}{dk_T^2 dx_1 dx_2} = x_1 \bar{G}(x_1) x_2 G(x_2) \frac{d\hat{\sigma}}{dk_T^2 dx_1 dx_2}$$

where $d\hat{\sigma}$ is evaluated at the parton level. Setting up the k_T integrals we have

$$\frac{d\sigma}{dx_1 dx_2} = \int d^2 k_T x_1 G(x_1) x_2 G(x_2) \frac{d\hat{\sigma}}{dk_T^2 dx_1 dx_2} \left(1 - \frac{5k_T^2 L_A}{2E(1 - x_1)} + \dots \right)$$

The second terms contains the nuclear effects; we have inserted the bound $x_g = k_T^2 L_A / 2E$. Since the integrand is proportional to k_T^2 , we can do the integral to estimate the A -dependent correction as

$$\frac{1}{A} \frac{d\sigma_A}{dx_1 dx_2} \sim \frac{d\sigma_1}{dx_1 dx_2} \left(1 - \frac{5 \langle k_T^2 \rangle_A L_A}{2E(1 - x_1)} + \dots \right) .$$

The correction has a size set by $\langle k_T^2 \rangle_A$. For nuclei $L_A \sim 1.2 \text{ Fm } A^{1/3}$, and the data for Υ_{1s} production gives $\langle k_T^2 \rangle_A \sim 0.16 A^{1/3} \text{ GeV}^2 + \langle k_T^2 \rangle_D$. The kinematics of producing quarkonium with invariant mass Q^2 and momentum fraction x_F requires $x_F = x_1 - x_2$, $Q^2/s = x_1 x_2$, which in the limit $Q^2/s \ll 1$ gives $x_F \approx x_1$. Then for the estimate for Υ_{1s} production

$$\frac{1}{A} \frac{d\sigma_A/dx_F}{d\sigma_1/dx_F} \approx 1 - \frac{5 \langle k_T^2 \rangle_A L_A}{2E(1-x_F)} .$$

This crude estimate does surprisingly well. Take for example $A = 184$ for tungsten, and $x_F = 0.5$. Then the effects of including energy loss lead to about 60% of the events compared to a calculation neglecting energy loss.

To check the analytic estimate we did some numerical calculations. We wish to compare the experimental trends in the data's x -dependence with its k_T dependence. We will use the experimental value of $\langle k_T^2 \rangle_A - \langle k_T^2 \rangle_{\text{int}}$, where $\langle k_T^2 \rangle_{\text{int}}$ is the intrinsic value, to calculate Δx_1 . This can then be used to calculate numerically the shift in the gluon distribution functions due to energy loss, thereby yielding the x dependence in the nuclear medium. Thus we take from the data

$$\Delta x_1 \leq L_A (\langle k_T^2 \rangle_A - \langle k_T^2 \rangle_{\text{int}}) / (2E_p) . \quad (18)$$

where we set $\langle k_T^2 \rangle_{\text{int}}$ to be equal to 0.91 GeV^2 in analogy to Drell Yan [18].

The x -integrated k_T distribution of the quarkonia has been parametrized experimentally [1] by,

$$f_A(k_T^2) = \xi(A) \left[\frac{1}{1 + (k_T/p_0)^2} \right]^6 \quad (19)$$

where $\xi(A)$ is an A dependent normalization factor and p_0 is also A dependent. The average value of p_T^2 , defined by,

$$\langle k_T^2 \rangle_A = \frac{\int_0^\infty dk_T^2 k_T^2 f_A(k_T^2)}{\int_0^\infty dk_T^2 f_A(k_T^2)} , \quad (20)$$

is equal to $p_0^2/4$. As given in Ref. [1], the values of p_0 for ${}^2\text{H}$ are

$$\begin{aligned} p_0 &= 2.78 \text{ for Drell Yan ,} \\ p_0 &= 3.22 \text{ for } \Upsilon . \end{aligned}$$

As discussed in the introduction, the values of p_0 are unfortunately not available for charmonium. The details of the p_T^2 dependence are not our object here, and we would prefer to insert them from data. To proceed with charmonium we will assume that it has the same transverse momentum dependence as bottomonium-an assumption which can be relaxed when data is obtained. The value of $\langle k_T^2 \rangle_A$ for the case of bottomonium increases as $0.16A^{1/3}$ [14]. The experimental fit to this data for the case of Υ is given in eq. (3).

We next need the production cross section in terms of the parton distributions. We assume that gluons give the most important contribution to quarkonium production [19] for moderately large x_f - this is discussed in detail later. The cross section integrated over transverse momentum is then given by

$$\frac{d\sigma}{dQ^2 dx_F} = \frac{x_1 x_2}{x_1 + x_2} G(x_1 + \Delta x_1) G(x_2) \frac{\sigma^{gg \rightarrow c\bar{c}}(Q^2)}{Q^2}. \quad (21)$$

We apply the same procedure to the Drell Yan continuum dilepton data, substituting quark and anti-quark distributions for the gluon distributions. The transverse momentum is quite different, but has been re-fit to match the data. The overall normalization is not relevant because we report ratios of nuclear targets to the proton. This procedure, repeated for each nucleus, gives us definite predictions for the x - dependence of the experimental data, which is discussed in the next section.

5. RESULTS

In this section we discuss the results of our simple GM energy loss calculation combined with the modified BH rule. The x_F dependence of the ratio $(d\sigma_A/dx_F)/d\sigma_D/dx_F$ extracted by using the experimentally measured k_T dependence is shown in Figs. (4-7) for the cases of Drell Yan, Υ_{1s} and charmonium respectively. For the bound we set Δx_1 (Eq. 18) to its maximum allowed value. We note that the experimentally measured points are well within this theoretical limit; as mentioned earlier, the trend in the data is to run parallel to the bound. For the case of charmonium we have taken the transverse momentum distribution to be the same as for bottomonium. We see that this choice fits the charmonium x_F dependence and therefore

gives the minimum value of the transverse momentum for the case of charmonium. The numerical results also include shadowing besides energy loss. We assumed the following functional form for the ratio $R_{\text{shadowing}}$ of nuclear to deuteron quark distributions due to shadowing,

$$R_{\text{shadowing}} = 0.809 + 0.261 \exp(-x_2 - 0.00526 A^{1/3}/x_2)$$

which fits the structure function data [20] with $\chi^2/\text{(degree of freedom)} = 0.86$. The shadowing for the case of gluons is included by assuming that it is the same as for quarks. The dashed lines in fig. (4), (5) and (6) represent the results without including energy loss. We see that although the Drell-Yan case is well reproduced by shadowing, charmonium and bottomonium data cannot be explained purely by shadowing, but require the addition of gluon energy loss.

Our results lead to several experimentally checkable predictions:

Small k_T

First, the conventional A^α analysis can be examined bin by bin in k_T and x_F . Generally speaking, the energy loss picture is distinguished by producing the largest suppression in the largest k_T regions. This does not mean that small k_T is totally safe, because rescattering can feed particles back in to this region. However we expect this to be controllable and therefore predict suppression increasing monotonically with k_T at fixed x_F or integrated over x_F . This implies that the suppression should be reduced if we consider the bin with k_T less than about 1 GeV 2 , where the dominant contribution comes purely from the primordial transverse momentum. Our estimate of the x_F distribution for different k_T cutoffs is given in Fig. (8). The perturbative part of transverse momentum is again calculated by subtracting the intrinsic contribution from the observed transverse momentum. At low transverse momentum, $k_T^2 < 2 < k_T^2 \rangle_{\text{int}}$ we calculated the limiting value of Δx_1 by setting k_T^2 equal to the intrinsic value of 0.91 GeV 2 .

The resulting curves, in Figure (8), show the ratios of the cross section as a function of x_F for various values of k_T . The curves are on a log plot, because the overall normalization is not being predicted. A shift in normalization N ,

translating the plots up or down, can be considered a free parameter. Our object is the shape of the curves, which clearly evolves with the k_T cut. For each region the experimental data should lie above the corresponding curve and follow the trend indicated in Fig. (8).

Large x_F

The idea of dominance of gluons in heavy quarkonia production is hardly new but has been raised again recently by BQV [7]. They consider the lowest order perturbative subprocess cross sections for producing a quark-antiquark pair with invariant mass Q^2 :

quark-antiquark channel:

$$\hat{\sigma}^{q\bar{q}}(Q^2) = \frac{2}{9} \frac{4\pi\alpha_s^2}{3Q^2} \left(1 + \frac{1}{2}\gamma\right) \sqrt{1 - \gamma},$$

gluon channel:

$$\hat{\sigma}^{gg}(Q^2) = \frac{\pi\alpha_s^2}{3Q^2} \left[\left(1 + \gamma + \frac{1}{16}\gamma^2\right) \log\left(\frac{1 + \sqrt{1 - \gamma}}{1 - \sqrt{1 - \gamma}}\right) - \left(\frac{7}{4} + \frac{31}{16}\gamma\right) \sqrt{1 - \gamma} \right],$$

where $\gamma = 4m_c^2/Q^2$. Convoluting these cross sections with standard parton distributions, BQV claim that the quark initiated process dominates over the glue-glue one for $x_F \gtrsim 0.5$. This conclusion is based on the perturbative normalizations given above, and the fact that the quark distributions fall less rapidly with x than the gluons. Quarks are forced kinematically to dominate as x_F goes to 1.

We agree with this in principle, but disagree that the crossover point can be given by the Born term calculation. Long experience with detailed calculations of quarkonium production at high energies favors gluons over quarks, indicating phenomenologically that the perturbative normalizations are not to be trusted too literally. As already noted, the data does not allow the option of ascribing the quarkonium suppression to quark channels while simultaneously accommodating the dilepton continuum. Moreover, the Born term does not even give the Drell Yan cross section correctly; for a long time it has been known that a “K-factor” of about 2 is needed to fit the data. Current understanding of K-factors is that they summarize higher order corrections from initial and or final state interactions. One cannot

assume the K-factors cancel out in nuclear ratios: what is relevant is the relative amount of quark and gluon contribution in each target.

There are several ways to proceed. One can estimate the crossover between quark versus gluon dominated quarkonium production with the K-factor method. Using Ref. [21], in the initial state interaction between two gluons we find a K-factor which is bigger than the annihilation K-factor by $N_c/C_f = 9/4$. With this (crude) estimate, the crossover point for quark annihilation channels over gluon channels at 800 GeV can be estimated to be 0.65. One may also treat this as an adjustable parameter to be determined experimentally once more data becomes available.

The value of large x_F is that one can experimentally “tune” the production process to favor quark initiated reactions. We have already noted that the Drell-Yan data indicates almost negligible energy loss, and smaller transverse momentum, for light quarks compared to gluons. If this is correct, then as x_F is increased above the crossover point the suppression in the nuclear medium should diminish. In Fig (9) we present a calculation illustrating the effect. This calculation was performed by using the Eichten *et al.* [22] parametrization of the quark and gluon distributions. If the final state effects are negligible than the data should show a sudden change in the current trend and an enhancement at large x_F . This is a dramatic signal meriting a careful search. For this calculation, we used the Born term cross sections modified by a relative normalization of 9/4 for the gluons to the quarks.

Our approach has combined theory with empirical patterns taken from the data. One could ask why light quarks do not come close to saturating the energy loss bound while apparently gluons do. The answer is, we don’t pretend to know. In the same vein, one can ask whether final state heavy quark energy losses should have been included. The answer is, the data does not indicate that significant energy loss from the heavy quarks needs to be introduced. Nevertheless, toward developing a truly model independent procedure, let us observe that the limit $x_F \rightarrow 1$ plays a key role. Suppose the up-turn as $x_F \rightarrow 1$ predicted above does not occur, even at such large values of x_F that we know the production is quark dominated. Then the suppression of heavy quarkonium compared to Drell–Yan production *must* be due to the heavy quark interactions above. Similarly, comparing experiments

with different beams – and especially pion beams known to be richer in quarks as $x \longrightarrow 1$ – has the potential to help separate the final state from the initial state effects. The Badier *et al.* data on pion initiated reactions [2] is consistent with this trend. We recommend using the data itself to separate the issues in a systematic way.

Certainly the effects discussed here have a direct bearing on the use of charmonium or quarkonium suppression as a probe of quark- gluon plasma formation at RHIC. Certainly a more thorough theoretical understanding and further experimental investigation of the phenomena is required before firm conclusions could be drawn from quarkonium production in heavy ion collisions.

Note Added: After this work was completed we became aware of a recent paper by M. S. Kowitt et al., Phys. Rev. Lett. **72**, 1318 (1994), which has extended the experimental data on x_F to larger values than was available previously. Their results show that the ratio of nuclear to Deuterium production starts to rise up considerably beyond $x_F = 0.65$ but then falls again around $x_F = 0.9$. Except for the point at $x_F = 0.9$, this data seems to support our picture. Comparing the data to our fig. (9), this suggests the idea that both light and heavy quarks lose negligible energy compared to gluons may be correct.

Acknowledgements: This work was supported in part under Department of Energy Grant Nos. DE-FG02-85ER-40214 and DE-FG05-91ER-40636 and the *Kansas Institute for Theoretical and Computational Science*.

Figure Captions

Fig. 1 Energy loss by emission of a parton in the presence of a source, marked by a circled “x”.

Fig. 2 Space-time picture of the null plane uncertainty principle $\Delta x^+ \Delta q^- \geq 1$
The light-front time interval Δx^+ is bounded to be less than about $\sqrt{2}L_A$ for ultra-relativistic processes.

Fig. 3 Energy loss in a theory with a fundamental 4-point vertex; the momentum fractions are indicated by the lengths of the lines. The uncertainty principle bounds the formation rate of the fastest forming parton (formation rate $k_T^2/2x_1P$) but says nothing further about the loss occurring among the other two partons, provided the sum of their formation rates $\Gamma_i = k_T^2/2x_iP$ is smaller than the first.

Fig. 4a-d Ratio of the nuclear to Deuterium cross section for Drell-Yan continuum dileptons (DY) as a function of x_f calculated by including the contribution only due to shadowing (dashed curve) and due to shadowing and the maximum allowed value of energy loss (solid curve) for A=12, 40, 56 and 184, respectively.

Fig. 5 The parameter α for bottomonium as a function of x_f calculated by including the contribution only due to shadowing (dashed curve) and due to shadowing and the maximum allowed value of energy loss (solid curve).

Fig. 6a-d Ratio of the nuclear to Deuterium cross section for J/ψ production as a function of x_f calculated by including the contribution only due to shadowing (dashed curve) and due to shadowing and the maximum allowed value of energy loss (solid curve) for A=12, 40, 56 and 184, respectively, using the transverse momentum distribution observed in the case of bottomonium.

Fig. 7 Predictions for the limiting values of ratios of nuclear to Deuterium cross sections for Υ_{15} production as a function of x_F . The data should lie above the curves.

Fig. 8 Limiting values of the ratio of Tungsten to Deuterium cross section for J/ψ production as a function of x_f for different transverse momentum bins. The slopes of the curves at small transverse momentum are much smaller than the slopes at higher transverse momentum. Data in each transverse momentum bin should have a slope less than or equal to that of the plotted curve. The overall x_f independent normalization factor N is not predicted.

Fig. 9 Ratio of Tungsten to Deuterium cross section for J/ψ production as a function of x_f including contributions of gluon fusion and quark-antiquark annihilation channels. The quark-antiquark annihilation contribution overtakes charmonium production for $x_f > 0.65$, producing the upturn in the curves.

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